

## BPS D-branes after Tachyon Condensation

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Toyonaka, Osaka, 560-0043, Japan***Abstract**

We construct an effective action describing brane-antibrane system containing  $N$  D-branes and  $N$   $\bar{D}$ -branes. BPS equations for remaining D-branes after tachyon condensation are derived and their properties are investigated. The value of the D-brane tension and the number of brane bound states are discussed.

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## 1. Introduction

The importance of tachyon in string theory has been shown in recent researches. The presence of the tachyon indicates the instability of brane-antibranes systems [1] and this system decays into a vacuum [2]. This is explained that the tension of the brane-antibranes system is cancelled by the minimum value of the tachyon potential  $V(T)$ , and the occurrence of such cancellation is examined by using string field theory [3]. The study of the tachyon condensation also provides new objects in string theory; non-BPS D-branes [2]. To classify these objects as well as the ordinary BPS D-branes K-theory is introduced into string theory and has been intensively studied [4].

In our previous paper [5], we discussed BPS D-branes after tachyon condensation in terms of a spacetime effective theory. D-branes were described as classical solutions of some BPS equations, and such description could deal with both parallel branes and intersecting branes. In addition, the form of the tachyon potential  $V(T)$  was determined by requiring the stability of the classical solutions.

We will generalize in this paper the effective theory of the brane-antibranes system to the one which describes the system containing  $N$  Dp-branes and  $N$   $\bar{\text{Dp}}$ -branes. This is a non-Abelian gauge theory coupled to the tachyon field, and this theory captures several properties of the brane system which cannot be represented by our old one.

This paper is organized as follows. We briefly review our previous work in section 2. The action of the multiple brane-antibranes system is constructed and the supersymmetry of this theory is discussed in section 3. The properties of the solutions of BPS equations and the fluctuations around them are investigated in section 4, and the asymptotic form of the tachyon field is determined in section 5. Section 6 is devoted to discussions.

## 2. BPS branes after tachyon condensation

The brane-antibranes system has shown to be unstable. The instability is due to the tachyonic mode coming from the string stretched between brane and antibranes. The calculations using string field theory show that the tension of the system is completely cancelled after tachyon condensation [2]. This means that the brane-antibranes system will decay into some stable state. The resulting stable configuration is either a vacuum or lower dimensional branes, and the latter can be classified by K-theory [4].

In our previous paper [5], we discussed the BPS D-branes which remain after tachyon condensation by using a spacetime action. The action of D9- $\bar{\text{D9}}$  system is the following,

$$S = \frac{1}{g^2} \int d^{10}x \left\{ -\frac{1}{4} F_{\mu\nu}^{(r)} F^{(r)\mu\nu} + \frac{i}{2} \bar{\psi}^{(r)} \Gamma^\mu \partial_\mu \psi^{(r)} - |D_\mu T|^2 - V(T) \right\}, \quad (1)$$

where  $D_\mu T = \partial_\mu T - iA_\mu^{(r)} T$ .  $A_\mu^{(r)}$  and  $\psi^{(r)}$  are ten-dimensional  $U(1)$  gauge multiplet on the brane-antibranes system (so called relative  $U(1)$ ), and  $T$  is the tachyon field. There also exists a charged massless fermion which is ignored here because it is irrelevant for

the following discussions. The action of Dp- $\bar{\text{Dp}}$  system can be obtained from (1) by dimensional reduction.

According to the arguments in string theory, this system should contain a stable solution which preserves some of the supersymmetries although the action (1) is not supersymmetric. Now let us consider the following “supersymmetry” transformations,

$$\delta A_\mu^{(r)} = \frac{i}{2} \bar{\epsilon} \Gamma_\mu \psi \quad (2)$$

$$\delta \psi^{(r)} = -\frac{1}{4} F_{\mu\nu}^{(r)} \Gamma^{\mu\nu} \epsilon - \frac{1}{2} (|T|^2 - \zeta) \Gamma^1 \Gamma^2 \epsilon \quad (3)$$

$$\delta T = 0 \quad (4)$$

where  $\zeta > 0$ . The action  $S$  is not invariant under the above transformations, and one of the conditions for  $\delta S = 0$  is

$$D_1 T + i D_2 T = 0. \quad (5)$$

$\delta \psi = 0$  for general  $\epsilon$  means

$$F_{12}^{(r)} + |T|^2 - \zeta = 0. \quad (6)$$

Thus we conclude that the solutions of (5) and (6) are the BPS solutions preserving 16 supercharges. (5)(6) are the Nielsen-Olesen vortex equations [6] and the solutions can be interpreted as BPS D-branes of codimension two.

### 3. Non-Abelian generalization

In this section we will construct an effective action of the brane-antibrane system which consists of  $N$  D9-branes and  $N$   $\bar{\text{D9}}$ -branes. This system is described by a gauge theory with gauge group  $SU(N) \times SU(N) \times U(1)$ . The relevant fields on the branes are the followings:  $A_\mu^{(1)}, \psi^{(1)}$  are  $SU(N)$  gauge multiplet on D9-branes and  $A_\mu^{(2)}, \psi^{(2)}$  are  $SU(N)$  gauge multiplet on  $\bar{\text{D9}}$ -branes.  $A_\mu^{(r)}, \psi^{(r)}$  are the relative  $U(1)$  multiplet.  $T$  is the tachyon filed in  $(\mathbf{N}, \bar{\mathbf{N}})$  representation of  $SU(N) \times SU(N)$  and this has the unit  $U(1)$  charge.

The action of this system is

$$\begin{aligned} S = \frac{1}{g^2} \int d^{10}x & \left[ -\frac{1}{4} \text{tr} F_{\mu\nu}^{(1)} F^{(1)\mu\nu} + \frac{i}{2} \text{tr} \bar{\psi}^{(1)} \Gamma^\mu D_\mu \psi^{(1)} - \frac{1}{4} \text{tr} F_{\mu\nu}^{(2)} F^{(2)\mu\nu} \right. \\ & \left. + \frac{i}{2} \text{tr} \bar{\psi}^{(2)} \Gamma^\mu D_\mu \psi^{(2)} - \frac{1}{4} F_{\mu\nu}^{(r)} F^{(r)\mu\nu} + \frac{i}{2} \bar{\psi}^{(r)} \Gamma^\mu \partial_\mu \psi^{(r)} - \text{tr} D_\mu T^\dagger D^\mu T - V(T) \right]. \end{aligned} \quad (7)$$

The field strengths and the covariant derivatives are defined as follows.

$$F_{\mu\nu}^{(\alpha)} = \partial_\mu A_\nu^{(\alpha)} - \partial_\nu A_\mu^{(\alpha)} - i[A_\mu^{(\alpha)}, A_\nu^{(\alpha)}] \quad (\alpha = 1, 2) \quad (8)$$

$$F_{\mu\nu}^{(r)} = \partial_\mu A_\nu^{(r)} - \partial_\nu A_\mu^{(r)} \quad (9)$$

$$D_\mu \psi^{(\alpha)} = \partial_\mu \psi^{(\alpha)} - i[A_\mu^{(\alpha)}, \psi^{(\alpha)}] \quad (10)$$

$$D_\mu T = \partial_\mu T - i A_\mu^{(1)} T + i T A_\mu^{(2)} - \frac{i}{\sqrt{N}} A_\mu^{(r)} T \quad (11)$$

The action (7) is completely determined by the gauge invariance and the Lorentz invariance except for the tachyon potential  $V(T)$ . As in the  $N = 1$  case, there exists a massless fermion in  $(\mathbf{N}, \bar{\mathbf{N}})$  representation of  $SU(N) \times SU(N)$  and the unit  $U(1)$  charge, which is also ignored as in our previous paper. We have used the following normalization of the generators  $t_a$  of  $SU(N)$ .

$$\text{tr } t_a t_b = \delta_{ab} \quad (12)$$

Then let us consider the “supersymmetry” transformations.

$$\delta A_\mu^{(\alpha)} = \frac{i}{2} \bar{\epsilon} \Gamma_\mu \psi^{(\alpha)} \quad (13)$$

$$\delta A_\mu^{(r)} = \frac{i}{2} \bar{\epsilon} \Gamma_\mu \psi^{(r)} \quad (14)$$

$$\delta \psi^{(1)} = -\frac{1}{4} F_{\mu\nu}^{(1)} \Gamma^{\mu\nu} \epsilon - \frac{1}{2} (T T^\dagger - \frac{1}{N} \text{tr} T T^\dagger \cdot 1) \Gamma^1 \Gamma^2 \epsilon \quad (15)$$

$$\delta \psi^{(2)} = -\frac{1}{4} F_{\mu\nu}^{(2)} \Gamma^{\mu\nu} \epsilon \quad (16)$$

$$\delta \psi^{(r)} = -\frac{1}{4} F_{\mu\nu}^{(r)} \Gamma^{\mu\nu} \epsilon - \frac{1}{2\sqrt{N}} (\text{tr} T T^\dagger - N\zeta) \Gamma^1 \Gamma^2 \epsilon \quad (17)$$

$$\delta T = 0 \quad (18)$$

The variation of the action (7) is

$$\begin{aligned} \delta S = \frac{1}{g^2} \int d^{10} x \frac{i}{2} \text{tr} [ & D_\mu T T^\dagger \bar{\psi}^{(1)} (i\Gamma^\mu - \Gamma^\mu \Gamma^1 \Gamma^2) \epsilon \\ & - i T^\dagger D_\mu T \bar{\psi}^{(2)} \Gamma^\mu \epsilon + \frac{1}{\sqrt{N}} D_\mu T T^\dagger \bar{\psi}^{(r)} (i\Gamma^\mu - \Gamma^\mu \Gamma^1 \Gamma^2) \epsilon + h.c. ]. \end{aligned} \quad (19)$$

Therefore the conditions for  $\delta S = 0$  are

$$D_1 T + i D_2 T = 0 \quad (20)$$

$$D_i T = 0 \quad (i = 0, 3, \dots, 9) \quad (21)$$

$$T \psi^{(2)} = 0. \quad (22)$$

Under the conditions (20)-(22), the action  $S$  is invariant for general  $\epsilon$ . When we would like to discuss classical solutions which preserve only part of 16 supercharges, the above conditions are relaxed in a suitable way.

#### 4. Gauge symmetry on the D-branes

We will investigate in this section classical solutions which preserve 16 supercharges. In particular we will focus on the solutions corresponding to D-branes of codimension two. For such solutions we assume that they depend only on the coordinates  $x^1$  and  $x^2$ , and

$A_{1,2}^{(\alpha)}, A_{1,2}^{(r)}$  are the only nonzero components of the gauge fields. Then  $\delta\psi^{(\alpha)} = 0, \delta\psi^{(r)} = 0$  mean

$$F_{12}^{(1)} + TT^\dagger - \frac{1}{N} \text{tr} TT^\dagger \cdot 1 = 0 \quad (23)$$

$$F_{12}^{(2)} = 0 \quad (24)$$

$$F_{12}^{(r)} + \frac{1}{\sqrt{N}} (\text{tr} TT^\dagger - N\zeta) = 0. \quad (25)$$

We set  $A_\mu^{(2)} = 0$  and define the  $U(N)$  gauge fields

$$A_\mu = A_\mu^{(1)} + \frac{1}{\sqrt{N}} A_\mu^{(r)} \cdot 1. \quad (26)$$

Then the eqs.(23)(25) become

$$F_{12} + TT^\dagger - \zeta \cdot 1 = 0. \quad (27)$$

This is a natural generalization of the Nielsen-Olesen vortex equation (5)(6).

The energy of the solutions of eqs.(20)(27) can be calculated as follows.

$$\begin{aligned} E = \frac{1}{g^2} \int d^{10}x & \left[ \frac{1}{2} \text{tr} (F_{12} + TT^\dagger - \zeta \cdot 1)^2 + \text{tr} |D_1 T + iD_2 T|^2 + \zeta \text{tr} F_{12} \right. \\ & \left. + V(T) - \frac{1}{2} \text{tr} (TT^\dagger - \zeta \cdot 1)^2 \right] \end{aligned} \quad (28)$$

Thus the solutions are stable topologically if

$$V(T) = \frac{1}{2} \text{tr} (TT^\dagger - \zeta \cdot 1)^2 \quad (29)$$

Now we will consider the properties of the solutions of eqs. (20)(27). For simplicity, let  $A_{1,2}$  and  $T$  be diagonal matrices.

$$A_\mu = \text{diag}(A_{(1)\mu}, \dots, A_{(N)\mu}) \quad (30)$$

$$T = \text{diag}(T_{(1)}, \dots, T_{(N)}) \quad (31)$$

Then eqs.(20)(27) become  $N$  sets of the ordinary vortex equations (5)(6). In this case the energy is

$$E = \sum_{i=1}^N \frac{\zeta}{g^2} \int d^{10}x F_{(i)12}. \quad (32)$$

The position of the vortex with respect to  $F_{(i)12}$  is determined as the point where  $T_{(i)} = 0$ . Therefore the number  $n(x)$  of D-branes at some point  $x$  is

$$n(x) = N - \text{rank}(T(x)). \quad (33)$$

This means that the solutions of eqs.(20)(27) can describe coincident D-branes as well as separated D-branes.

Now we will discuss the fluctuations around the classical solutions. Suppose that  $T$  has the form

$$T = \text{diag}(0, \dots, 0, T_{n+1}, \dots, T_N) \quad (34)$$

at some point  $x$ . Consider the fluctuations of  $A_\mu^{(2)}, \psi^{(2)}, A_\mu^{(r)}, \psi^{(r)}$ ,

$$A_\mu^{(2)} = A_\mu^{(2)c} + a_\mu^{(2)} \quad (35)$$

$$A_\mu^{(r)} = A_\mu^{(r)c} + a_\mu^{(r)} \quad (36)$$

$$\psi^{(2)} = \varphi^{(2)} \quad (37)$$

$$\psi^{(r)} = \varphi^{(r)} \quad (38)$$

where the superscript  $c$  indicates the classical solution of eqs.(20)(27). For the dynamics of the fluctuations to be supersymmetric, the conditions (20)-(22) have to be satisfied, that is,

$$T(a_\mu^{(2)} + \frac{1}{\sqrt{N}}a_\mu^{(r)} \cdot 1) = 0 \quad (39)$$

$$T\varphi^{(2)} = 0. \quad (40)$$

Let us define the  $U(N)$  gauge multiplet

$$a_\mu = a_\mu^{(2)} + \frac{1}{\sqrt{N}}a_\mu^{(r)} \cdot 1 \quad (41)$$

$$\varphi = \varphi^{(2)} + \frac{1}{\sqrt{N}}\varphi \cdot 1. \quad (42)$$

Then the conditions (39)(40) imply

$$a_\mu = \begin{pmatrix} a_{n,\mu} & 0 \\ 0 & 0 \end{pmatrix}, \quad \varphi = \begin{pmatrix} \varphi_n & 0 \\ 0 & 0 \end{pmatrix}, \quad (43)$$

where  $a_{n,\mu}, \varphi_n$  are  $n \times n$  hermitian matrices. This indicates that  $U(n)$  gauge symmetry appears on the  $n$  coincident D-branes. In fact, they are the fields in the super Yang-Mills theory with the ordinary supersymmetry transformations. If  $T$  has the generic form, we have to set  $a_\mu = 0, \varphi = 0$  to satisfy the conditions (39)(40), thus the fluctuations are restricted on the branes.

We have not mentioned the fluctuations  $a_\mu^{(1)}$  of  $A_\mu^{(1)}$ . If we include them, the action contains a term  $\text{tr}[A_\mu^{(1)c}, a_\nu^{(1)}]^2$ . This indicates that in general  $a_\mu^{(1)}$  become massive and can be ignored. However it may be possible for them to become massless, for example, at the brane intersection, if we consider such brane systems.

## 5. Asymptotic behavior of the tachyon field

In this section we turn to the solutions which preserve part of 16 supercharges.

First we will consider the solutions corresponding to intersecting branes. For example, there is an intersecting brane configuration which preserves 8 supercharges; two D7-branes extended along (01256789) and (03456789). The remaining supersymmetries correspond to the parameter  $\epsilon$  satisfying

$$\Gamma^1 \Gamma^2 \epsilon = \Gamma^3 \Gamma^4 \epsilon. \quad (44)$$

In this case the conditions for  $\delta S = 0$  under the transformations (13)-(18) are modified as follows.

$$D_1 T + i D_2 T = 0 \quad (45)$$

$$D_3 T + i D_4 T = 0 \quad (46)$$

$$D_i T = 0 \quad (i = 0, 5, \dots, 9) \quad (47)$$

$$T \psi^{(2)} = 0 \quad (48)$$

We assume, as in the previous section, that  $A_\mu^{(2)} = 0, A_i^{(1)} = 0, A_i^{(r)} = 0$  and the solutions depend only on the coordinates  $x^1, x^2, x^3, x^4$ . Then  $\delta\psi^{(1)} = 0, \delta\psi^{(r)} = 0$  mean

$$F_{12}^{(1)} + F_{34}^{(1)} + T T^\dagger - \frac{1}{N} \text{tr} T T^\dagger = 0 \quad (49)$$

$$F_{13}^{(1)} - F_{24}^{(1)} = 0 \quad (50)$$

$$F_{14}^{(1)} + F_{23}^{(1)} = 0 \quad (51)$$

$$F_{12}^{(r)} + F_{34}^{(1)} + \frac{1}{\sqrt{N}} (\text{tr} T T^\dagger - N \zeta) = 0 \quad (52)$$

$$F_{13}^{(r)} - F_{24}^{(r)} = 0 \quad (53)$$

$$F_{14}^{(r)} + F_{23}^{(r)} = 0. \quad (54)$$

In terms of the  $U(N)$  gauge fields (26) the above equations can be rewritten as

$$F_{12} + F_{34} + T T^\dagger - \zeta \cdot 1 = 0 \quad (55)$$

$$F_{13} - F_{24} = 0 \quad (56)$$

$$F_{14} + F_{23} = 0. \quad (57)$$

Eqs.(45)(46)(55)-(57) are the equations which describe the intersecting branes. One can show that the solutions of them are stable topologically when one takes  $V(T) = \frac{1}{2} \text{tr}(T T^\dagger - \zeta \cdot 1)^2$ .

We introduce the complex coordinates.

$$z^k = x^{2k-1} + i x^{2k} \quad (k = 1, 2) \quad (58)$$

Then the BPS equations become as follows,

$$D_{\bar{a}}T = 0 \quad (59)$$

$$F_{\bar{a}\bar{b}} = 0 \quad (60)$$

$$-ig^{a\bar{b}}F_{a\bar{b}} + TT^\dagger - \zeta \cdot 1 = 0, \quad (61)$$

where  $g_{a\bar{b}}$  is the metric in the complex coordinates; its nonzero components are  $g_{z^1\bar{z}^1} = g_{z^2\bar{z}^2} = \frac{1}{2}$ .

Similar equations can be derived for the solutions which preserve only 4 supercharges. Consider an intersecting brane configuration which contains D7-branes extended along (01234789), (03456789) and (01256789). The conditions for the remaining supersymmetries in this case are

$$\Gamma^1\Gamma^2\epsilon = \Gamma^3\Gamma^4\epsilon = \Gamma^5\Gamma^6\epsilon. \quad (62)$$

The BPS equations are formally the same as that for D7-D7' system except that in this case  $a, b$  take  $z^1, z^2, z^3$ .

Eqs.(59)(60) can be easily solved. From (60)  $A_{\bar{a}}$  are determined to be

$$A_{\bar{a}} = iV\partial_{\bar{a}}V^{-1}, \quad (63)$$

where  $V$  is an  $N \times N$  complex matrix. Let  $T = VT_0$ . Then (59) becomes

$$\partial_{\bar{a}}T_0 = 0. \quad (64)$$

Thus  $T_0$  is just a holomorphic matrix.

Before discussing eq.(61), we would like to fix the boundary condition at spatial infinity. It should be imposed that  $V(T) = 0$  at infinity, which implies that  $T$  becomes proportional to a unitary matrix. We take the polar decomposition  $V = UH$  and set  $T_0 = \sqrt{\zeta} \cdot 1$ , where  $U$  is unitary and  $H$  is hermitian. Then the boundary condition means that  $H \rightarrow 1$  at infinity.

We define  $R = V^\dagger V = H^2$ . Then eq.(61) can be rewritten in terms of  $R$ .

$$g^{a\bar{b}}\partial_{\bar{b}}(R^{-1}\partial_a R) = \zeta(R - 1) \quad (65)$$

The solutions of this equation in  $N = 1$  case are discussed in [5].  $U$  will be determined by requiring the regularity of the solutions.

To the intersecting brane configuration including three kinds of D7-branes considered above, we can add D3-branes extended along (0789) without breaking more supersymmetries. Therefore eqs.(59)-(61) are expected to have the corresponding solutions and they will have nonzero third Chern character. From the boundary condition,

$$A_\mu \rightarrow iU\partial_\mu U^\dagger. \quad (66)$$

The third Chern character can be calculated by the following expression.

$$\int_{\mathbf{R}^6} \text{tr} F^3 = -\frac{i}{10} \int_{\partial\mathbf{R}^6} \text{tr} U dU^\dagger dU dU^\dagger dU dU^\dagger \quad (67)$$



Let us discuss the special case  $N = 4$ , in which  $U$  can be explicitly constructed,

$$U = \frac{1}{r} \gamma_m x_m \quad (m = 1, \dots, 6), \quad (68)$$

where  $\gamma_m$  are the lower-left parts of the  $SO(6)$  gamma matrices,

$$\Gamma_m = \begin{pmatrix} & \gamma_m^\dagger \\ \gamma_m & \end{pmatrix} \quad (69)$$

and  $r^2 = (x^1)^2 + \dots + (x^6)^2$ . This gives the asymptotic form of the gauge fields whose third Chern character is integrated to be unity, corresponding to the single D3-brane on the intersecting D7-branes.

The asymptotic behavior of the tachyon field is also determined in this case as follows.

$$T \rightarrow \frac{\sqrt{\zeta}}{r} \gamma_m x_m \quad (70)$$

This is what has been discussed in the literature [4].

## 6. Discussions

In the previous section, we derived the equations (59)-(61) which described intersecting brane systems. Thus, by quantizing the collective coordinates of the solutions, the number of their bound states may be able to be determined, as will be explained below, if we assume their existence. In fact solving the BPS equations is not an easy task. However the number can be deduced from a few properties of the solutions and supersymmetry.

The BPS equations are not scale invariant. Therefore all of the bosonic moduli of the solutions will correspond to the positions of them and the scale of them will be fixed. Since there remains some part of supersymmetry, the number  $F$  of fermionic partners of the bosonic moduli will be a multiple of the number  $B$  of the bosonic moduli. The low energy dynamics of the brane system will be described by quantum mechanics whose target space is the moduli space of the solutions, and in particular, the number of bound states will be determined by counting the number of ground states of quantum mechanics.

Consider first D7-D7' system discussed in the previous section. In this case, it is expected that  $B = 4$  and  $F = 8$ . This implies that the number of the bound states is  $2^{\frac{8}{2}} = 16$ . Since D7-D7' system is a T-dual of D4-D0 system, this result is appropriate [7]. Next consider D7-D7'-D7''-D3 system. In this case  $B = 6$ ,  $F = 6$  and the number of the bound states is  $2^{\frac{6}{2}} = 8$ . This is the value expected from the counting of a black hole entropy [8] because this system is a T-dual of D4-D4'-D4''-D0 system.

In our discussions, the tachyon potential  $V(T)$  is determined by requiring the stability of the brane solutions. Thus we can calculate the tensions of the remaining branes as well as that of the original brane-antibrane system, and check whether our model reproduces

the correct values. It is natural to set  $g^{-2} = (2\pi\alpha')^2\tau_9$  and  $\zeta = \frac{1}{2\alpha'}$ , where  $\tau_9$  is the tension of the BPS D9-brane. Then the tension of the original brane-antibrane system is

$$V(0) = \frac{\zeta^2}{2g^2}Tr(1) = \frac{\pi^2}{4} \cdot N \cdot 2\tau_9, \quad (71)$$

and the tension of the resulting codimension-two brane is

$$\frac{2\pi\zeta}{g^2} = \pi\tau_7. \quad (72)$$

These do not match the correct values. Probably this is because we have considered the action (7), which includes only the leading terms in order of  $\alpha'$ . Therefore if one can construct an effective action containing all  $\alpha'$  corrections, then we will be able to determine the form of the tachyon potential  $V(T)$  by the same way as explained in this paper, and the model will provide the correct values for the D-brane tensions. The string field theory calculation [3] seems to imply that the string loop corrections are not important, and our classical arguments will be enough to do the job.

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